

The μ and soft masses from the intermediate scale brane with non-factorizable geometry

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Abstract

The proton decay problem and the negative brane tension problem in the original Randall-Sundrum model can be resolved by interpreting the Planck scale brane as the visible sector brane. The hierarchy problem is resolved with supersymmetry, and the TeV scales for soft masses and μ in supersymmetric models are generated by the physics at the intermediate scale ($\sim 10^{11-13}$ GeV) brane.

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One of the major theoretical puzzles in particle physics is the gauge hierarchy problem [1], which is basically the Higgs boson mass parameter (TeV scale) problem in the presence of the fundamental scale of order the Planck mass ($M = 2.44 \times 10^{18}$ GeV). The well-known technicolor and supersymmetric solutions aim toward obtaining a TeV scale scalar mass naturally [2]. Recently, an alternative solution toward the gauge hierarchy problem has been suggested by Randall and Sundrum (RS) [3], where the huge gap between the Planck and TeV scales is explained by the exponential warp factor of a 5 dimensional space-time metric,

$$ds^2 = e^{-2\sigma(y)} \eta_{\mu\nu} dx^\mu dx^\nu + b_0^2 dy^2, \quad (1)$$

where y is the fifth dimension coordinate and $\sigma(y) = kb_0|y|$ is given as a solution of the Einstein equation. For the purpose of obtaining such a metric, Randall and Sundrum assumed a negative bulk cosmological constant $\Lambda_b \equiv -6k^2M^3$, and required a S^1/Z_2 symmetry in the extradimension y . Then the metric becomes non-factorizable. In this model, there are two branes (orbifold fixed points) as 4 dimensional boundaries, Brane 1 (B1) with a positive cosmological constant $\Lambda_1 \equiv 6k_1M^3$ at $y = 0$, and Brane 2 (B2) with a negative cosmological constant $\Lambda_2 \equiv 6k_2M^3$ at $y = y_c$. Although the RS setup introduces cosmological constants, their solution is still static because of the fine-tuning between the bulk and brane cosmological constants, $k = k_1 = -k_2$, which is a consistency condition in the model.

On the branes, there could exist ‘brane fields’ that live only in the concerned brane. They correspond to the twisted sector fields in string theory, which are necessary for anomaly freedom of the theory after orbifold compactification [4]. Brane fields on B2 are then governed by $M_{Pl}e^{-\sigma(y=y_c)} \approx$ TeV scale physics, while bulk fields like the graviton and brane fields on B1 are governed by Planck scale physics [3].

Even though the two scales are easily understood under the RS setup, it still has a few problems:

(i) Late cosmology demands that the visible brane is better to have a positive cosmological constant [5], which corresponds to B1 in the RS model. Since it is inconsistent with

the RS's original motivation that the TeV brane(B2) is the visible brane, there have been several propositions [6–8] such that B2 has a positive cosmological constant.

(ii) Under the RS setup it is difficult to rescue the GUT idea because the cutoff scale at B2 is TeV. It is well-known that the GUT scale should be around 3×10^{16} GeV with the particle content of the minimal supersymmetric standard model(MSSM).

(iii) In addition, the RS model has the proton stability problem since the relevant interaction scale is expected to be TeV. For $\tau_p > 10^{33}$ years, one has to forbid proton decay operators up to dimension 14 [6,7], which is a difficult problem even though it may be achievable in contrived models.

To circumvent the above difficulties within the non-factorizable geometry, we propose to regard *the Planck brane (B1) as the visible brane*, in which the standard model (SM) particles live, and obtain *the TeV scale parameter(s) from a source at B2*. To obtain the TeV scale mass parameter(s) naturally, it would be desirable to *forbid the required TeV mass parameter(s) at B1*. The role of B2 is to generate the source for TeV scale masses of B1. The effects of B2 is transmitted to B1 by bulk fields.

For this purpose, it would be necessary to introduce an additional symmetry that forbid unwanted mass parameter(s) at B1, and a bulk field as a messenger, which may couple to both brane fields. In our simple model, we will assume a global $U(1)_A$ symmetry and enforce the messenger bulk field to carry the $U(1)_A$ quantum number. This symmetry will be broken spontaneously due to an interaction with brane fields at B2 which plays the role of symmetry breaking source. Then the messenger bulk field will get vacuum expectation value(VEV) to give the Higgs a TeV scale mass parameter.

Even if we get a TeV scale at tree level, in order to guarantee the stability between the Planck and TeV scales, we introduce supersymmetry(SUSY). Then, the scale problem in the theory becomes the μ problem, which is the scale problem in supergravity (SUGRA) models [9]. With $\mu = 0$, there exists the Peccei-Quinn(PQ) symmetry [10]. Therefore, the electroweak symmetry breaking would introduce the unwanted Peccei-Quinn-Weinberg-Wilczek axion [11]. With the μ term the PQ symmetry is explicitly broken, and there does

not appear the unwanted axion in the model. In our case, however, a TeV scale axion decay constant could appear if we generate the μ term at a TeV scale brane. Therefore, we do not introduce a TeV scale brane at all. Instead, we introduce an intermediate scale brane B2. This intermediate scale is of order 10^{11-13} GeV. The TeV scale at B1 is generated by the bulk field(s) coupled to Higgs doublets *through non-renormalizable interaction* in the model described below.

The low energy SUGRA models also have another scale, the soft mass parameters which are expected to be around TeV scale. These soft masses are generated once SUSY is broken. The popular SUGRA models use the ideas of hidden sector gaugino condensation at intermediate scale, gauge mediated SUSY breaking, or SUSY breaking due to anomaly. In our case, we employ the gaugino condensation *at B2*, which seems to be the simplest method. At B2, the natural mass scales are around the intermediate scale due to the warp factor. If the β function of a nonabelian gauge group is large and negative, this gauge group confines immediately below the cutoff scale of B2 and the corresponding gauginos condense around the intermediate mass scale. In string inspired $E_8 \times E'_8$ models, the extra factor group E'_8 is suitable for this purpose. In supergravity, gravitino mass is a barometer for the strength of SUSY breaking. In our case, the gravitino mass would be of order TeV scale,

$$m_{3/2} \approx \frac{\langle \lambda \lambda \rangle}{M_P^2} \approx TeV , \quad (2)$$

for the gaugino condensation scale of order 10^{13} GeV. The soft parameters are also of order TeV scale since *the gravitino is a bulk field*. We will discuss it again later. Thus, we can show that TeV scale SUSY breaking as well as TeV scale μ term can be realized easily in the RS setup *even if we identify B1 as the visible brane*.

In this paper, we neglect the backreaction by the brane and bulk fields to the background geometry, and we do not discuss the cosmological constant problem, which is assumed to be zero by the fine-tuning in the RS model. The cosmological constants on B1 and B2 are non-zeros. However, we will ignore them also because there exists a massless mode of the graviton, which gives the *flat space action* effectively after integrating over the extradimension [7,12].

This situation does not change in the supersymmetric generalization of the RS model [13] the method of which is used in this paper.¹

Let us suppose the global supersymmetry for a simple analysis and introduce three siglet fields with a global symmetry $U(1)_A$ which plays the role of the PQ symmetry. The $U(1)_A$ charges, A , of the various fields are shown in Table 1.

	Bulk	B1		B2	
Fields	Σ^i	H_1, H_2	MSSM	S	Z
A	+1	-1	$\frac{1}{2}$	-1	0

Table 1. The $U(1)_A$ charges, A of the various fields in the bulk and branes.

All the MSSM particles are B1 brane fields, Σ^i live(s) in the bulk and S and Z are B2 brane fields. Of course, the gravity multiplet which is not shown in Table 1 is of course the bulk field. Brane fields $S \equiv (\phi_S, \psi_S)$ and $Z \equiv (\phi_Z, \psi_Z)$ form $D = 4$ chiral superfields at B2. Bulk fields $\Sigma^i \equiv (\Phi^i, \Psi_{(L,R)})$ constitute $D = 5$ hypermultiplet, where Φ^i ($i = 1, 2$) are two complex scalars and $\Psi_{(L,R)}$ are left and right handed Dirac fermions. We define $\Psi_{L,R}$ as $\gamma_5 \Psi_{L,R} = \pm \Psi_{L,R}$. They constitute a $N = 2$ supermultiplet and make the theory non-chiral.

At B1, all dimension 2 and 3 operators containing $H_1 H_2$, including the so-called μ term, are forbidden by the $U(1)_A$ symmetry. The dominant term in the superpotential consistent with the $U(1)_A$ symmetry is the dimension 4 operator,

$$\sim \frac{\Sigma^2}{M_P} H_1 H_2 . \quad (3)$$

This is just a *schematic formula*; it should be rewritten such that *the interactions between the bulk and brane fields respect $D = 4$ SUSY after integrating over the extradimension*, which is described below. At B2, the most general dimension 3 superpotential is

$$\sim Z \left(\Sigma S - M_P^2 \right) , \quad (4)$$

¹ Alternative supersymmetrizations of the RS model in the framework of 5D supergravity are suggested in [14].

which is also a *schematic formula*.

The action of the hypermultiplet in the bulk is

$$S_{bulk} = - \sum_i \int d^5x \sqrt{-G} \left[g^{MN} \partial_M \Phi^{i*} \partial_N \Phi^i + \frac{i}{2} \left(\bar{\Psi} \Gamma^M \nabla_M \Psi - (\nabla_M \bar{\Psi}) \Gamma^M \Psi \right) \right. \\ \left. + M_{\Phi^i}^2 |\Phi^i|^2 + M_{\Psi} \bar{\Psi}_L \Psi_R + M_{\Psi} \bar{\Psi}_R \Psi_L \right] , \quad (5)$$

where $\Gamma^M \equiv e_a^M \gamma^a$. ∇_M is the covariant derivative on a curved manifold and defined as $\nabla_M \equiv \partial_M + \omega_M$, where ω_M is the spin connection,

$$\omega_\mu = \frac{1}{2} \gamma_5 \gamma_\mu \frac{d\sigma(y)}{dy} \quad \text{and} \quad \omega_5 = 0 \quad . \quad (6)$$

Invariance under the supersymmetry transformations [17,18],

$$\delta \Phi^i = i\sqrt{2} \epsilon^{ij} \bar{\eta}^j \Psi \quad (7)$$

$$\delta \Psi = \sqrt{2} \left[\Gamma^M \partial_M \Phi^i \epsilon^{ij} - \frac{3}{2} \sigma' \Phi^i (\epsilon \sigma_3)^{ij} - M_{\Psi} \Phi^i \epsilon^{ij} \right] \eta^j , \quad (8)$$

requires that the five dimensional masses of the scalars and fermions satisfy

$$M_{\Phi^1}^2 = (t^2 + t - \frac{15}{4}) \sigma'^2 + (\frac{3}{2} - t) \sigma'' \quad (9)$$

$$M_{\Phi^2}^2 = (t^2 - t - \frac{15}{4}) \sigma'^2 + (\frac{3}{2} + t) \sigma'' \quad (10)$$

$$M_{\Psi} = t \sigma' , \quad (11)$$

where $\sigma' \equiv kb_0 \left[2 \left(\theta(y) - \theta(y - y_c) \right) - 1 \right]$ and $\sigma'' \equiv 2kb_0 \left(\delta(y) - \delta(y - y_c) \right)$, and t is an arbitrary dimensionless parameter. In flat space-time, supersymmetry requires the same masses for the scalars and fermions. But in the AdS_5 background, the fields in the same supermultiplet must have different masses [17,18].

Equations of motion for the above action allow the massless modes [18],

$$\Phi^{1,(0)}(x, y) = \frac{e^{(3/2-t)\sigma(y)}}{\sqrt{2b_0 y_c N}} \phi_{\Sigma}(x) \quad (12)$$

$$\Psi_L^{(0)}(x, y) = \frac{e^{(2-t)\sigma(y)}}{\sqrt{2b_0 y_c N}} \psi_{\Sigma}(x) , \quad (13)$$

where $\sigma(y) \equiv kb_0 |y|$ and the normalization factor N is given by

$$N^2 \equiv \frac{e^{\sigma_c(1-2t)} - 1}{\sigma_c(1-2t)} , \quad (14)$$

where $\sigma_c \equiv kb_0y_c$ and N is reduced to 1 when $t = 1/2$. Note that *the massless modes depend on 'y'*. $\Phi^2(x, y)$ and $\Psi_R(x, y)$ do not have massless modes since they are inconsistent with the orbifold condition [18], and hence the fermion mass terms cannot exist for the massless modes. We will see below that the scalar mass terms are canceled also in the effective 4 dimensional action. When decoupling the Kaluza-Klein massive modes, $N = 2$ SUSY breaks down to $N = 1$ SUSY for the massless modes of the hypermultiplet, and $\Sigma^{1,(0)} \equiv (\Phi^{1,(0)}(x, y), \Psi_L^{(0)}(x, y))$ form $N = 1$ supersymmetric chiral multiplets.

Inserting Eqs. (9)–(14) to Eq. (5) and integrating by parts, we get the 4 dimensional effective action for the massless modes,

$$\begin{aligned} S_{bulk}^{eff} &= - \int d^4x \int_{-y_c}^{y_c} dy \left[\frac{b_0 e^{-4\sigma}}{2b_0 y_c N^2} \left(e^{2\sigma} e^{(3-2t)\sigma} \eta^{\mu\nu} \partial_\mu \phi_\Sigma^* \partial_\nu \phi_\Sigma + e^{(4-2t)\sigma} e^\sigma i \bar{\psi}_\Sigma \gamma^\mu \partial_\mu \psi_\Sigma \right) \right. \\ &\quad \left. + \frac{b_0}{2b_0 y_c N^2} \left(-e^{(\frac{3}{2}-t)\sigma} \partial_y (e^{-4\sigma} \partial_y e^{(\frac{3}{2}-t)\sigma}) |\phi_\Sigma|^2 + e^{-4\sigma} M_{\Phi^1}^2 e^{(3-2t)\sigma} |\phi_\Sigma|^2 \right) \right] \\ &= - \left(\frac{1}{2y_c N^2} \int_{-y_c}^{y_c} e^{(1-2t)\sigma(y)} dy \right) \int d^4x \left[\eta^{\mu\nu} \partial_\mu \phi_\Sigma^* \partial_\nu \phi_\Sigma + i \bar{\psi}_\Sigma \gamma^\mu \partial_\mu \psi_\Sigma \right] \\ &= - \int d^4x \left[\eta^{\mu\nu} \partial_\mu \phi_\Sigma^* \partial_\nu \phi_\Sigma + i \bar{\psi}_\Sigma \gamma^\mu \partial_\mu \psi_\Sigma \right] , \end{aligned} \quad (15)$$

where we see the scalar mass terms are eliminated and the contributions of $\partial_y \Psi$ and $\partial_y \bar{\Psi}$ add up to zero. Thus, we confirm that the ϕ_Σ and ψ_Σ are massless fields. In the above equations, t is not fixed yet. However, if the bulk fields are required to couple *supersymmetrically* to the brane fields on a brane, t should be fixed to $\frac{1}{2}$ and the bulk fields form a 4 dimensional supermultiplet as

$$\bar{\Sigma}(x, y) = \left(e^{\sigma(y)} \phi_\Sigma(x), e^{\frac{3}{2}\sigma(y)} \psi_\Sigma(x) \right) , \quad (16)$$

which will become clear below.

At the intermediate brane B2, the brane fields are required to be rescaled such that their kinetic terms have the canonical forms,

$$S_{B2}^{kin} = - \sum_{i=S,Z} \int d^4x e^{-4\sigma_c} \left[e^{2\sigma_c} \eta^{\mu\nu} \partial_\mu \phi_i^* \partial_\nu \phi_i + e^{\sigma_c} i \bar{\psi}_i \gamma^\mu \partial_\mu \psi_i \right]$$

$$= - \sum_{i=S,Z} \int d^4x \left[\eta^{\mu\nu} \partial_\mu \tilde{\phi}_i^* \partial_\nu \tilde{\phi}_i + i \bar{\tilde{\psi}}_i \gamma^\mu \partial_\mu \tilde{\psi}_i \right] , \quad (17)$$

where we use $e^a{}_\mu|_{y=y_c} = e^{-\sigma_c} \delta_\mu^a$ and $\tilde{\phi}_i$ and $\tilde{\psi}_i$ are defined as the rescaled fields as follows,

$$\phi_i \equiv e^{\sigma_c} \tilde{\phi}_i \quad \text{and} \quad \psi_i \equiv e^{\frac{3}{2}\sigma_c} \tilde{\psi}_i . \quad (18)$$

At B2, the brane fields $S = (\phi_S, \psi_S)$ and $Z = (\phi_Z, \psi_Z)$, and bulk fields $\bar{\Sigma}^{1,(0)}|_{y=y_c} = (e^{\sigma_c} \phi_\Sigma(x), e^{\frac{3}{2}\sigma_c} \psi_\Sigma(x))$ can form Yukawa interaction terms *supersymmetrically* in the action,

$$\begin{aligned} S_{B2}^{int} &= \int d^4x \sqrt{-g_4} \left[((e^{\sigma_c} \phi_\Sigma) \psi_S \psi_Z + \phi_S \psi_Z (e^{\frac{3}{2}\sigma_c} \psi_\Sigma) + \phi_Z (e^{\frac{3}{2}\sigma_c} \psi_\Sigma) \psi_S + \text{h.c.}) \right. \\ &\quad \left. - |\phi_S \phi_Z|^2 - |\phi_Z (e^{\sigma_c} \phi_\Sigma)|^2 - |(e^{\sigma_c} \phi_\Sigma) \phi_S|^2 + M_P^2 ((e^{\sigma_c} \phi_\Sigma) \phi_S + \text{h.c.}) - M_P^4 \right] \\ &= \int d^4x \left[(\phi_\Sigma \tilde{\psi}_S \tilde{\psi}_Z + \tilde{\phi}_S \tilde{\psi}_Z \psi_\Sigma + \tilde{\phi}_Z \psi_\Sigma \tilde{\psi}_S + \text{h.c.}) \right. \\ &\quad \left. - |\tilde{\phi}_S \tilde{\phi}_Z|^2 - |\tilde{\phi}_Z \phi_\Sigma|^2 - |\phi_\Sigma \tilde{\phi}_S - m_I^2|^2 \right] , \end{aligned} \quad (19)$$

where we use $\sqrt{-g_4}|_{y=y_c} = e^{-4\sigma_c}$ and Eq. (18). m_I is defined as $m_I \equiv M_P e^{-\sigma_c} \sim 10^{11} - 10^{13}$ GeV for $\sigma_c = k b_0 y_c \approx 11.5 - 16$. Here we use the Weyl spinor notation for spinor fields. Eq. (19) is clearly invariant under the $D = 4$ SUSY transformations. To see that explicitly, let us define

$$\tilde{\Sigma}(x) \equiv \left(\phi_\Sigma(x), \psi_\Sigma(x) \right), \quad \tilde{S}(x) \equiv \left(\tilde{\phi}_S(x), \tilde{\psi}_S(x) \right) \quad \text{and} \quad \tilde{Z}(x) \equiv \left(\tilde{\phi}_Z(x), \tilde{\psi}_Z(x) \right). \quad (20)$$

Then we can derive a superpotential at B2 from Eq. (19) as follows,

$$W_{B2} = \tilde{Z} \left(\tilde{\Sigma} \tilde{S} - m_I^2 \right) . \quad (21)$$

Here we note that if we did not choose $t = \frac{1}{2}$ in Eqs. (12) and (13), we could not get the $D = 4$ effective supersymmetric interactions between bulk and brane fields.

At B1 all the needed Yukawa interactions in the MSSM except the μ term are allowed. The nonrenormalizable interactions are expected in the supergravity generalization of our global SUSY study. In this case, we may define the superpotential W as the maximum [19] allowed by the symmetry

$$G = K + M_P^2 \log \frac{|W|^2}{M_P^6} = \text{invariant} \quad (22)$$

where K is a Kähler potential. In this case, we generally introduce all possible non-renormalizable terms consistent with the symmetry. Typically, the following superpotential is present,

$$W_{B1} = \frac{\tilde{\Sigma}^2}{M_P} H_1 H_2 . \quad (23)$$

From Eqs. (21) and (23), we derive supersymmetric scalar potential,

$$\begin{aligned} V = V_{B1} + V_{B2} &\simeq \sum_{p=\Sigma, H_1, H_2} \left| \frac{\partial W_{B1}}{\partial \phi_p} \right|^2 + \sum_{i=\Sigma, \tilde{S}, \tilde{Z}} \left| \frac{\partial W_{B2}}{\partial \phi_i} \right|^2 \\ &= \left| \frac{\phi_\Sigma^2}{M_P} \phi_{H_1} \right|^2 + \left| \frac{\phi_\Sigma^2}{M_P} \phi_{H_2} \right|^2 + \left| \frac{2\phi_\Sigma}{M_P} \phi_{H_1} \phi_{H_2} \right|^2 + |\tilde{\phi}_S \tilde{\phi}_Z|^2 + |\tilde{\phi}_Z \phi_\Sigma|^2 + |\phi_\Sigma \tilde{\phi}_S - m_I^2|^2 . \end{aligned} \quad (24)$$

The potential has a minimum at $\phi_\Sigma \tilde{\phi}_S = m_I^2$ and $\tilde{\phi}_Z = \phi_{H_1} = \phi_{H_2} = 0$. But TeV scale SUSY breaking would lift this degeneracy as much as TeV scale. Thus ϕ_Z , ϕ_{H_1} and ϕ_{H_2} could have VEV of TeV scale, and the electroweak mass scales are derived. The PQ symmetry breaking scale is at the intermediate scale, and there results a very light axion [20].

Since ϕ_Σ obtains an intermediate scale VEV ², we can get a TeV scale μ term at B1,

$$W_{B1} \approx \frac{\langle \phi_\Sigma \rangle^2}{M_P} H_1 H_2 . \quad (25)$$

This form of the μ term from the symmetry principle has been considered before [9,21,19]. Here, we realize it with the intermediate scale brane world.

Similarly, we expect a TeV scale gravitino mass. Let us suppose that the gravitino mass is generated at B2. At the Planck brane B1 a bilinear scalar field $\phi^* \phi$ couples to two gravitino lines with a coupling suppressed by the Planck mass M_{Pl} . These gravitinos propagate in the bulk and at B2 they meet to produce the gravitino mass. Since the bulk propagator cutoff ranges up to M_{Pl} , the resulting effective gravitino mass after the Feynman integration would be of order TeV scale since the gravitino mass at B2 is of order TeV scale. But without

²In a model such that the μ term is generated from $W_{B1} = \frac{\langle \phi_\Sigma \rangle^3}{M_P^2} H_1 H_2$, 10^{13} GeV would be preferred as an intermediate scale.

the gravitino, the soft mass would not be generated. For example, the graviton coupling to $\phi^*\phi$ at $B1$ is through a derivative of the scalar field, which means that for the vanishing momentum of ϕ the bulk propagation of the graviton would give a vanishing potential and hence no soft mass contribution. Namely, the SUGRA is needed to obtain the soft mass terms of the scalars living at $B1$. For the gravitino propagation calculation, however, we need an exact expression for the gravitino propagator in the bulk. Even though we have that expression, we must cutoff at the Planck scale for the above gravitino loop and the estimate is anyway a kind of order of magnitude.

Therefore, it is better and clear to use an effective 4 dimensional theory after integrating the y coordinate. Here, we will use the SUGRA language. In order to trigger the SUSY breaking with the gaugino condensation, let us introduce on $B2$ a vector multiplet (corresponding to the confining hidden sector force), $(A_\mu^a(x), \chi^a \equiv e^{\frac{3}{2}\sigma_c} \lambda^a(x))$. We can easily check that $A_\mu^a(x)$ and $\lambda^a(x)$ give the canonical kinetic terms through the above procedure, Eq. (17)–(18).

The graviton and gravitino, which are bulk fields, have the massless modes ($e^a{}_\mu(x, y) \sim e^{-\sigma(y)} \tilde{e}^a{}_\mu(x)$; $\Psi_\mu^{(1,2)}(x, y) \sim e^{-\frac{1}{2}\sigma(y)} \psi_\mu(x)$) [13,15,12,7]. Thus, after decoupling the KK modes and integrating over the extra dimension and following the above procedure Eq. (5)–(15), we can confirm that the $e^a{}_\mu$ and ψ_μ are the effective 4 dimensional graviton and gravitino, which are actually massless at least without any other physics [13],

$$\begin{aligned} S_{bulk}^{kin} &= \int d^4x \int_{-y_c}^{y_c} dy \left[\sqrt{-g} \frac{1}{2} R - \Lambda_b + \left(\frac{-\Lambda_b}{6} \right) \epsilon^{MNOPQ} \bar{\Psi}_O \Sigma_{PQ} D_Q \Psi_N + \dots \right] \\ &= \int d^4x \left[\sqrt{-\tilde{g}_4} \frac{M_{Pl}^2}{2} \tilde{R}_4 + \epsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu \bar{\sigma}_\nu D_\rho \psi_\sigma \right], \end{aligned} \quad (26)$$

where we set M as 1 and use the RS fine-tuning conditions between the bulk and brane cosmological constants, $\sqrt{-\Lambda_b/6} = \Lambda_1/6 = -\Lambda_2/6$. We note here that *the effective 4 dimensional space-time is flat*.

After decoupling the KK modes, however, the massless mode of gravitino and the *brane* gaugino field could compose the 4-fermion interaction terms at $B2$ in the supersymmetric way [16,15],

$$\begin{aligned}
S_{B2}^{4-fermi} &\sim \int d^4x \sqrt{-g_4} \left[\frac{e^{3\sigma_c} \langle \lambda^a \lambda^a \rangle}{M_P^2} e^{-\frac{1}{2}\sigma_c} \bar{\psi}_\mu \Sigma^{\mu\nu} e^{-\frac{1}{2}\sigma_c} \psi_\nu + \dots \right] \\
&= \int d^4x \left[\frac{\langle \lambda^a \lambda^a \rangle}{M_P^2} \bar{\psi}_\mu \sigma^{\mu\nu} \psi_\nu + \dots \right]
\end{aligned} \tag{27}$$

where we use $e^a{}_\mu|_{y=y_c} = e^{-\sigma_c} \delta_\mu^a$ and $\Sigma^\mu \equiv e^\mu{}_a \sigma^a = (e^{\sigma_c} \delta_a^\mu) \sigma^a$. Thus the gravitino mass term, which could be a barometer for SUSY breaking, is generated in our effective 4 dimensional theory when the brane gauginos condense. Since the intermediate scale ($\sim 10^{13} \text{GeV}$) is a natural cutoff and the condensation scale on B2 *with a large gauge group*, the above term gives a TeV scale gravitino mass successfully, as explained before.

After integrating over the extra dimension, the hidden brane loses its geometrical meaning and just remains as a *hidden sector gauge field*. Thus we obtain an effective 4 dimensional theory again *with a visible gauge sector whose typical cutoff scale is the Planck scale, and a hidden gauge sector whose cutoff scale is the intermediate scale*. The 4 dimensional gravity sector fields, which was in the bulk before, interact with all the fields irrespective of which sector they come from. Therefore, in the effective theory language it is *exactly the same picture with the conventional D=4, N=1 SUGRA scenario with the hidden sector, except that the hidden sector's cutoff scale is the intermediate scale*. Thus, the soft mass terms of the scalar fields are the square of the gravitino mass $m_{3/2}$ which is the same as the B_2 gravitino mass given in Eq. (27) since the y integration for $m_{3/2}$ gets a contribution only from B_2 through the Dirac delta function $\delta(y - y_c)$.

The effective 4 dimensional theory gives possibly a consistent D=4, N=1 SUGRA. In this case, all kinds of SUSY breaking coefficients in the Lagrangian are parameterized with the gravitino mass. Therefore, we could get the TeV scale SUSY breaking effects from the TeV scale gravitino mass through the gravitational interaction in the bulk.

In conclusion, we constructed a successful brane model where the visible sector lives in the Planck scale brane, and the supersymmetry breaking occurs at the intermediate scale brane. A TeV brane is not needed. The TeV scale, the soft masses and μ , are derived quantities from the mass source at the intermediate scale brane, through the mediation by bulk fields. The proton decay problem is resolved since the grand unification is achieved at

the Planck brane. Since the visible sector is at B1, the problem of negative brane tension in the original RS model does not arise. For the stability of soft masses and μ , we introduced supersymmetry, which is possible for a specific value of the parameter t .

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